CERN-TH.95/134
OUTP-95-15P
hep-ph/9506283
(revised)

Successful Supersymmetric Inflation

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Abstract

We reconsider the problems of cosmological inflation in effective supergravity theories. A singlet field in a hidden sector is demonstrated to yield an acceptable inflationary potential, without fine tuning. In the simplest such model, the requirement of generating the microwave background anisotropy measured by COBE fixes the inflationary scale to be about 10^{14} GeV, implying a reheat temperature of order 10^5 GeV. This is low enough to solve the gravitino problem but high enough to allow baryogenesis after inflation. Such consistency requires that the generation of gravitational waves be negligible and that the spectrum of scalar density perturbations depart significantly from scale-invariance, thus improving the fit to large-scale structure in an universe dominated by cold dark matter. We also consider the problems associated with gravitino production through inflaton decay and with other weakly coupled fields such as the moduli encountered in (compactified) string theories.

Accepted for publication in *Nuclear Physics B*

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1 Introduction

Although inflation [1] is an attractive solution to the horizon/flatness problems of the standard Big Bang model and to the cosmological monopole problem of GUTs, it has yet to find a compelling physical basis [2]. Interest in this question has been rekindled by the COBE [3] discovery of temperature fluctuations in the cosmic microwave background (CMB) consistent with a 'Harrison-Zeldovich' scale-invariant power spectrum. This arises naturally in 'slow-roll' inflationary models from quantum fluctuations of the scalar field which drives the de Sitter phase of exponential expansion, as it evolves towards its minimum [4]. The observed small amplitude, $\delta T/T \sim 10^{-5}$, requires an extremely flat scalar potential stabilized against radiative corrections. This picks out a gauge singlet field in a theory incorporating supersymmetry (SUSY) as the most likely candidate for the 'inflaton' [5, 6, 7].

However such models contain very weakly coupled fields having masses of $\mathcal{O}(m_W)$ and this creates difficulties with the cosmological history *after* inflation. For example, gravitinos can have observable effects on the standard cosmology since they decay very late with lifetime $\sim\ M_{\rm Pl}^2/m_{3/2}^3$ [8]. Although their primordial abundance can be inflated away, they are recreated during 'reheating' as the inflaton oscillates about its minimum, converting vacuum energy into radiation [9]. This imposes a severe constraint on the reheat temperature since even a small number of massive late decaying particles can disrupt primordial nucleosynthesis or the thermalization of the CMB [10]. Also inflation offers no obvious solution to the 'Polonyi problem' [11, 12] which is rather more subtle, being associated with the production of the unwanted states via vacuum decay in sectors whose vacuum energy is affected through gravitational interactions by the vacuum energy driving inflation. This is of particular relevance to the moduli in (compactified) string theories [13].

In this paper we introduce and discuss several novel features of inflation in superstring theories with moduli fields. The central question is whether acceptable inflation can occur in a natural way with no need for fine tuning or for multiple correlated periods of inflation. We argue that this is indeed the expectation in the presence of moduli fields. First we consider the possible scale for an inflationary potential in theories with a single stage of dynamical symmetry breakdown responsible for SUSY breaking in a 'hidden sector'. Contrary to commonly expressed opinions, it appears quite plausible for the scale of the inflationary potential to be of $\mathcal{O}(10^{14})$ GeV, consistent with the value needed to generate the observed CMB fluctuations. We then discuss whether fine tuning is neccessary in order to achieve a sufficiently flat potential. The novel feature considered here is the possibility that the initial conditions of moduli fields selected with just the values needed to generate successful inflation. This happens because the moduli fields are expected to have chaotic initial values, hence there will be *some* domains with the right initial conditions to give an inflationary potential. Such domains will inflate leading, through quantum fluctuations, to 'eternal' inflation; these configurations will thus dominate the final state of the universe. With this motivation we consider the resultant inflationary potential, using a simple example as a guide, and find that it has a significant curvature, at about the maximum possible for slow roll inflation.

We then consider the reheating phase and, in particular, the post-inflation production of gravitinos or other weakly coupled states. Demanding that their thermal production

during reheating be under control implies a general constraint on the inflationary potential and leads to the observationally testable prediction that there should be a negligible tensor component in the CMB anisotropy in models with a single stage of inflation. A second prediction following from the curvature of the inflationary potential is that the spectrum of scalar density perturbations should deviate from scale-invariance so as to suppress small-scale power. This provides a better fit to the observed clustering and motions of galaxies in a cold dark matter (CDM) universe. We demonstrate how even direct gravitino production is acceptably small in the favoured class of inflationary models in which the inflaton resides in a hidden sector of the theory. Finally we discuss the implications of the Polonyi problem, in particular for recent models which treat moduli as dynamic variables at the electroweak scale.

2 Supersymmetric Inflation

If there is a stage of inflation in the early universe driven by approximately constant vacuum energy, then requiring that the associated quantum fluctuations should not create anisotropies in the CMB in excess of those observed by COBE bounds the scale of this potential energy, $V^{1/4}$, to be at least two orders of magnitude below the Planck scale.[3] At this energy scale, gravitational interactions and interactions due to string and Kaluza-Klein states with masses of $\mathcal{O}(M_{\rm Pl})$ are small, hence one can use an effective field theory to describe the inflation sector. This effective theory should of course contain the Standard $SU(3)_{\rm c}\otimes SU(2)_{\rm L}\otimes U(1)_Y$ Model. The condition that the electroweak symmetry breaking scale should naturally be small ($\ll M_{\rm Pl}$) without fine tuning the parameters of the theory then requires it to be supersymmetric with SUSY broken (softly) only just above the electroweak scale. We shall consider in particular supergravity (SUGRA) theories based on local supersymmetry which include a description of gravity and are of the type that descend from the superstring. The $N=1$ SUGRA theory describing the interaction of gauge singlet superfields is specified by the Kähler potential $K(\Phi,\Phi^\dagger)$, which determines the form of the kinetic term in the Lagrangian. The scalar potential is [15]

$$V=\frac{1}{4}{\rm e}^K\left[G_a(K^{-1})^a_bG^b-3\mid W\mid^2\right]\ , \tag{1}$$

where the Kähler function $G_a=K_aW+W_a$, the indices $a,\ b$ denote derivatives with respect to the chiral superfields, Φ, and $W(\Phi)$ is the superpotential which specifies the Yukawa couplings of the theory and the scalar couplings related by supersymmetry.

2.1 "Natural" Inflation

A criticism that is often levelled at models of vacuum energy driven inflation is that they require unnatural fine-tuning [2]. There are two parts to this problem. First of all, such inflationary models necessarily contain small parameters in order to generate

[3]The generation of gravitational waves becomes increasingly pronounced as the inflationary energy scale approaches the Planck scale [4]. Assuming these to be entirely responsible for the observed quadrupole anisotropy yields the conservative bound $V^{1/4}\lesssim 4\times10^{-3}M_{\rm Pl}$ [14]. If the quadrupole is in fact due to scalar density fluctuations then the bound is even more restrictive as we shall see.

an inflationary potential significantly lower than the Planck scale. To provide a physical justification for this, one should perhaps associate the inflationary scale with one of the scales needed in viable unified field theories of the fundamental interactions. The second aspect of the problem concerns the initial conditions, specifically the question of how probable are the initial field configurations necessary to ensure an inflationary era. In this section we discuss both these issues in the context of supersymmetric inflation.

2.1.1 The inflationary scale

Let us first consider the mass scale(s) to be expected in the effective field theory which may play a role in setting the energy scale for the potential in the inflation sector. We know that in string and/or compactified theories there are numerous states with mass of $\mathcal{O}(M)$, where we use as our basic mass unit $M \equiv M_{\mathrm{Pl}}/\sqrt{8\pi} \simeq 2.44 \times 10^{18}\,\mathrm{GeV}$. In addition, the success of the supersymmetric unification prediction relating the strong and electroweak gauge couplings suggests there is a stage of spontaneous gauge symmetry breaking at a scale $M_X \approx 2 \times 10^{16}\,\mathrm{GeV}$ [16]. The theory must also have a source of SUSY breaking characterized by the gravitino mass, $m_{3/2} \lesssim 1\,\mathrm{TeV}$. The most plausible origin for the latter is dynamical supersymmetry breaking via nonperturbative effects driven by a new interaction which becomes strong at a scale Λ_{c}, thus inducing gaugino condensation in a hidden sector. In this case $m_{3/2} \sim \langle\lambda\lambda\rangle/M^2$ and the gaugino condensate $\langle\lambda\lambda\rangle \approx (10^{13}\,\mathrm{GeV})^3$. How could such scales affect a gauge singlet sector? Consider first the gauge symmetry breaking scale. We denote by $\chi, \bar{\chi}$, the gauge non-singlet fields which acquire a vacuum expectation value (vev) of $\mathcal{O}(M_X)$ along a D-flat direction thus breaking the gauge symmetry, and by Θ, $\bar{\Theta}$, the gauge singlet fields with masses of $\mathcal{O}(M)$. Allowing for a coupling between these fields, consider a superpotential of the form (suppressing coupling constants of order unity)

$$W = M\Theta\bar{\Theta} - \bar{\Theta}\chi\bar{\chi}\ , \tag{2}$$

from which we see that the gauge symmetry breaking vev in $\chi, \bar{\chi}$ induces a vev in the massive field given by $\langle\Theta\rangle \equiv \Delta = \langle\chi\rangle\langle\bar{\chi}\rangle/M \approx 10^{14}\,\mathrm{GeV}$. If Θ couples to other fields in the theory, they will acquire mass determined by this vev, the value of the mass determined by the strength of coupling.

The point is that once a stage of symmetry breaking is generated below the Planck scale in some sector of the theory, one may expect it to generate masses in every sector of the theory. However this conclusion may not be true during an inflationary era in a string theory. The reason is that in theories of this type all scales other than the Planck scale must arise dynamically. The most obvious way is through the supersymmetry breaking sector via non-perturbative effects, which become strong at a scale Λ_{c} far below the Planck scale, driving a gaugino condensate. However one must then require that all other related scales are consistent with the formation of the condensate. For example if a gauge symmetry breaking vev is generated it is due to a supersymmetry breaking trigger. A vev for the χ field introduced above may be induced through radiative breaking in which the χ soft-SUSY-breaking mass-squared term in the Lagrangian is driven negative by radiative corrections at some scale M_X and minimization of the effective potential for the $\chi, \bar{\chi}$ fields and the SUSY breaking sector together, requires the vev to be of $\mathcal{O}(M_X)$, essentially independent of the SUSY breaking trigger. However this only happens when

it is energetically favourable for the vev to develop and in particular *cannot* occur if there is a large potential energy associated with the inflation sector. In other words, the inflationary potential will cause χ to acquire a soft mass term inhibiting the development of its vev until inflation is over. These considerations place an *upper* limit on the scale of inflation, viz. the positive inflationary potential should not exceed the energy scale involved in gauge and SUSY breaking otherwise the latter will be inhibited by the very inflationary phase it is supposed to drive.

In the literature (e.g. ref.[17]) this constraint has been implicitly interpreted as requiring that the inflationary vacuum energy should not exceed the non-zero F- or D-term contribution to the effective potential which is responsible for SUSY breaking. This is given by $(m_{3/2}M_{\rm Pl})^2 \sim \Lambda_{\rm c}^6/M^2$, leading to an upper bound of $\mathcal{O}(10^{11})$ GeV for the scale of the inflationary potential, which is indeed too small. There are two possible ways out of this apparent impasse. It may be that the interaction in a second sector also becomes strong and dynamically generates a second, larger scale which is responsible for inflation but does not lead to SUSY breaking after inflation. We disfavour this possibility because we are trying to avoid multiple correlated sectors when constructing a "natural" theory. However it is quite reasonable to expect that a single scale of dynamical SUSY breaking has a *different* magnitude during inflation than after inflation. For example the inflationary potential may inhibit gauge symmetry breaking in the manner discussed above. If this symmetry breaking relates to the group which drives SUSY breaking (when its interaction becomes strong), the result is that the β-function of the group is bigger during inflation, hence the scale at which the coupling becomes strong is higher, in turn driving a larger scale of SUSY breaking.[4]

A second possibility is that the inflation scale is related to the scale, $\Lambda_{\rm c}$, at which the coupling becomes large and not to the final SUSY breaking scale. Indeed this seems to us the more reasonable choice for certainly $\Lambda_{\rm c}$ is a relevant scale of new physics, for example, connected with the appearance of confined massive states. Thus it is perfectly possible that due to the initial conditions the vacuum energy starts off at the natural binding energy scale $\Lambda_{\rm c}^4$ and only drops to the scale $\Lambda_{\rm c}^6/M^2$ as the fields (in this case the inflaton) adjust to their vacuum values. Clearly such a scenario requires a strong cancellation between the terms contributing to the vacuum energy at the minimum but this is just what is usually assumed about gaugino condensation. As we shall demonstrate, inflation at the $\mathcal{O}(10^{14})$ GeV scale is compatible with all constraints for a sub-class of inflationary models which are of the 'new' [18, 20] rather than of the 'chaotic' [19] type.[5] This is marginally consistent with the identification of the scale with $\Lambda_{\rm c} \sim 10^{13}$ GeV, given the uncertainties involved in this identification.

Within these constraints we consider it to be the expectation rather than the exception for there to be sectors in the theory associated with the mass scale required for acceptable

[4]Here we are assuming that the initial value of the gauge coupling set by the dilaton vev is fixed, say at a modular invariant point. There is, of course, the possibility that the dilaton vev also varies, its potential affected by the inflationary potential. This is another possible source for a change in the SUSY breaking scale before and after inflation.

[5]In order to realise the latter, one must rely on a small coupling constant (rather than a ratio of mass scales) to provide the scale of inflation in terms of the Planck scale. In string theories small couplings can indeed arise but only dynamically when moduli fields acquire large vevs. These however do not lead to inflationary potentials because the would be inflationary potential is not sufficiently flat in the moduli direction [21].

inflation. To demonstrate this we may use the mechanism discussed above to construct a model for inflation in which the scale of the potential is set by the original spontaneous breaking of the gauge symmetry (and, if this scale is not fundamental, by the associated SUSY breaking trigger, provided the scale is constrained in the manner just discussed). The starting point is the form of the potential describing the inflaton which, for the reasons discussed above, we take to be a gauge singlet, ϕ; at this stage we need not specify whether it is an elementary scalar field or arises as a composite object. The coupling between the superfield Φ containing the inflaton and Θ will be constrained by any symmetries of the theory. The superpotential of eq.(2) has an R-symmetry under which Θ and $\chi\bar{\chi}$ transform as $e^{i\gamma}$, $\bar{\Theta}$ transforms as $e^{i(2\beta-\gamma)}$ while the superspace coordinate transforms as $e^{-i\beta}$. Consider the sector with an R-singlet, gauge singlet, superfield Φ which contains the (complex) inflaton field, ϕ as its scalar component. The most general superpotential, P, describing Θ, $\bar{\Theta}$, χ and Φ, has the form dictated by this R-symmetry,

$$P = \Theta\bar{\Theta}M \ f\left(\frac{\Phi}{M}\right) + \bar{\Theta}\chi\bar{\chi} \ , \tag{3}$$

where $f(x)$ is a function which is not constrained by the R-symmetry alone. (We have absorbed the constant term generating the Θ mass in f.) As discussed above, once $\chi, \bar{\chi}$ acquire vevs breaking the gauge symmetry, the field Θ will also acquire a vev,

$$\Delta = \frac{\langle\chi\rangle\langle\bar{\chi}\rangle}{2M} \ , \tag{4}$$

(with $f(0) = 1$) leading to the inflaton superpotential

$$I\,(\Phi) = \Delta^2 M \ f\left(\frac{\Phi}{M}\right) . \tag{5}$$

Of course there are other ways to motivate this form; the most obvious is that it arises directly from the gaugino condensate sector in which the confined states have mass of $\mathcal{O}(\Delta)$ and the inflaton is one of these composite states. We shall consider a particular choice for $f(\Phi/M)$ in Section 3 but first we consider the general features of such a superpotential. The first point is that the associated potential for the inflaton ϕ is proportional to Δ^4, showing how the gauge breaking scale naturally sets the scale for the potential of light fields in the effective potential obtained by integrating out massive fields. By appealing to a richer symmetry structure one can construct examples in which the scalar potential is proportional to different powers of Δ, but this example suffices to illustrate the general point. Moreover it will turn out that this power is just what is required for acceptable inflation generating the observed amplitude of density perturbations.

2.1.2 Initial conditions

We turn now to the second part of the naturalness question, namely whether one requires fine tuning of the initial conditions or of the parameters determining the potential for there to be an inflationary era [22]. Consider the structure of the scalar potential following from the superpotential, I. It will be determined once the Kähler potential is specified but for now we need merely assume that it has at least one local maximum or point of inflection. If the initial conditions are such that the field lies very close to this point then there will

be an inflationary era. Of course if the initial value of the field is determined by the high temperature effective potential [23] then it is unlikely that this value coincides with the point at which the first derivative of the potential vanishes, unless one fine tunes the parameters of the potential to ensure this. If, however, thermal effects do not fix the initial state, then it will have, in general, a broad distribution [19, 22]. In this case, while there may only be a small probability that one starts with the value needed for an inflationary era, the region posessing this initial value *will* inflate and become the overwhelmingly probable state after inflation. This is the basic idea which leads us to the conclusion that weakly coupled fields which are not in thermal equilibrium below the Planck scale are overwhelmingly likely to generate inflation without the need for any fine tuning.

To quantify this, consider a field ϕ which drops out of thermal equilibrium at temperatures below the Planck scale. (Note that this is what we expect for the field with interactions described by the superpotential of eq.(5) because the leading cubic term giving a Yukawa coupling and associated quartic scalar couplings, has a coupling of $\mathcal{O}(\Delta^2/M^2)$.) At the Planck temperature it is reasonable to suppose that through gravitational effects the field has thermal fluctuations given by $\langle(\delta\phi)^2\rangle \sim M^2$ about some value set by the physics at the Planck scale.[6] Due to the spatial derivative terms in the Lagrangian the field will be smoothed out over a horizon volume by the Kibble mechanism. Thus as the universe cools below the Planck scale, one might suppose that the fields free stream and (averaging over the enlarged horizon) settle in to their Planck scale mean values with deviation $\propto M/\sqrt{n}$ at a temperature where the horizon contains n Planck volumes. However as the temperature drops the potential energy associated with the random initial conditions will become comparable to the thermal energy and the field expectation values will move to reduce the potential. This will be a complicated process, dependent on the initial field configurations, in which the various fields do *not* free stream but are correlated in any given horizon volume and as a result the fields are *not* likely to head to the mean values fixed at the Planck scale. Thus we conclude that the field values in different horizon volumes will, apart from sharing excluded regions of field space for which the potential energy is greater than the thermal energy, be largely uncorrelated until a (very low) temperature at which the excluded regions of field space grow sufficiently precisely to specify the allowed regions. The implication is that one should use a (flat) distribution of initial values, as at the Planck scale.

Let us now consider the case when the mean value of the field in any horizon volume lies at a local maximum or point of inflection of the potential, at the temperature for which the potential energy at this point is comparable to the thermal energy. This is just the situation discussed in ref. [24] wherein it was shown that despite the fluctuations of the field throughout the horizon volume it behaves as if it were a classical field with value given by the mean and evolving according to the equation of motion

$$\ddot{\phi} + 3H\dot{\phi} = -\frac{\partial V}{\partial \phi} \approx m_\phi^2 \phi \ , \tag{6}$$

[6] It has been argued [19] that a reasonable initial condition emerging from the physics at the Planck scale is a flat probability distribution for values of the field which have $V(\phi) \lesssim M^4$. For a generic potential $V = \lambda\phi^4$ with small coupling λ, the field value can be thus much larger than the Planck scale as is required in chaotic inflation. However in string theories couplings are of order unity so the initial field value should *not* be taken to be much larger than the Planck scale.

where we have expanded about the point of the potential for which the first derivative vanishes. Inflation will occur if the first term on the LHS is negligible (i.e. if $m_\phi < H$). Since the two quantities are typically of the same order, whether this happens or not depends on the detailed values of the parameters. Assuming that it does, inflation will continue until the curvature of the potential becomes significant at $\phi \sim M$. This determines the number of e-folds of inflation to be

$$N = Ht_{\text{inflation}} \sim \frac{3\Delta^4}{M^2 m_\phi^2} \ln\left(\frac{M}{\phi_0}\right), \tag{7}$$

where ϕ_0 is the initial value of ϕ. This is given by $\phi_0 \sim T_{\text{deS}}$ the temperature at which the thermal and false vacuum energy in the de Sitter phase are about equal [24]; thus $T_{\text{deS}} \sim \Delta$, which can be used in eq.(7) to determine N. If there are no very small parameters we have $m_\phi \sim H$, so inflation will be limited to a few tens of e-folds. Now we can complete the quantitative estimate of the probability that we end up with an inflationary universe. With our best guess for the initial conditions being an uniform distribution over the Planck scale, the probability we start with ϕ within a distance ϕ_0 of the maximum of the potential is just $\phi_0/M \sim \Delta/M \sim 10^{-4}$. However this is more than compensated for by an inflation of tens of e-folds, so the final state of the universe is dominated by states which have enjoyed an inflationary era.

As we have seen, the initial field configuration with vanishing derivative for the scalar potential dominates the final state of the universe due to the inflationary expansion associated with such a configuration. One may then consider the possibility that the second derivative also vanishes, for in this case the amount of inflation increases. In theories following from the superstring such a possibility appears to be quite likely because the couplings which determine the higher derivatives are themselves determined by the vevs of moduli fields. If the initial values of these vevs are random, there will be regions in which the initial conditions are such that the second derivative vanishes too. This process will not continue to an arbitrary degree for there are only a finite number of moduli and the condition that a higher derivative vanishes determines some combination of them. However it is quite plausible that there are sufficient moduli [7] undetermined by the time the inflationary potential dominates that at least the second derivative vanishes. (Then there will be many more e-folds of expansion than the $\sim 50-60$ needed for the observable universe to arise from one causally connected domain.) While it is possible that further derivatives may vanish, this is irrelevant in determining the implications for the observable universe, so in what follows we consider examples in which only the first two derivatives are zero.

In fact it has been pointed out to us (A. Linde, private communication) that this picture fits in very well with the concept of 'everlasting eternal inflation' [25] in which, during the slow-roll, quantum fluctuations continally produce regions *closer* to the maximum of the potential. Although the initial probability for this is small, these regions become exponentially larger than the remainder and thus ultimately dominate the final state of the universe. In this way the inflationary universe reproduces itself eternally, so essentially all final states descend from an inflationary era. In terms of the moduli, the

[7] In Calabi-Yau compactification there is one complex structure moduli for each (2,1) form and one Kähler moduli for each (1,1) form.

initial state is continually being driven to the optimal choice of moduli vevs for which the infationary potential is flattest. This is the ultimate "fine tuning without fine tuning"! We note that for *any* inflationary model it is essential to invoke a mechanism of this sort if one is to overcome the argument[26] that there is apparently a hopelessly small chance that the initial conditions of the universe are sufficiently homogeneous and isotropic for inflation to start at all.

2.2 Constraints on the inflationary potential

Before making a specific choice for $f(\Phi/M)$ in eq.(5) let us consider the overall constraints on the scale of inflation. In order to illustrate these, we specialize to the minimal choice of Kähler potential, $K = \Phi^\dagger\Phi$, corresponding to canonical kinetic energy for the scalar fields. (This is known to be the case in certain orbifold models, however here we make this choice for pedagogic reasons; the discussion can be extended to non-minimal Kähler potentials as well.) With this choice the scalar potential following from the superpotential, I, is [15]

$$V_I(\phi) = \mathrm{e}^{|\phi|^2/M^2}\left[\left|\frac{\partial I}{\partial\phi} + \frac{\phi^* I}{M^2}\right|^2 - \frac{3|I|^2}{M^2}\right]_{\Phi=\phi}. \tag{8}$$

In what follows we shall consider slow-roll models in which the function f can be expanded as a power series in ϕ/M or, as in simple chaotic inflation models where ϕ/M is large, is given by a combination of powers. If there are no small parameters other than Δ/M (as argued above, this is the simplest case and may eliminate the need for fine tuning) then V has a minimum at $\phi \sim M$. Clearly more complicated models can be constructed and we will comment on these later but our purpose here is to discuss whether the bounds on the reheat temperature imposed by consideration of the production and subsequent decay of unstable gravitinos present a problem for the generic model characterized by a *single* energy scale.

2.2.1 The gravitino problem

Gravitinos or other unwanted fields will be thermally excited during the reheating epoch following inflation and decay subsequently into high energy particles. The effects of such decays on cosmological observables such as the primordial abundances of helium and deuterium and the spectrum of the CMB have been studied in some detail. The most stringent constraint on the gravitino abundance comes from the observational upper limit on the abundances of D and ^{3}He which can be produced via photofission of ^{4}He by the radiation cascades from gravitino decay: $m_{3/2}(n_{3/2}/n_\gamma) \lesssim 3 \times 10^{-12}\,\mathrm{GeV}$ [27]. Taking the gravitino abundance produced during reheating by $2 \to 2$ processes involving gauge bosons and gauginos to be [9]

$$\frac{n_{3/2}}{n_\gamma} \simeq 2.4 \times 10^{-13}\left(\frac{T_\mathrm{R}}{10^9\,\mathrm{GeV}}\right), \tag{9}$$

this implies an upper bound on the reheat temperature of [27]

$$T_\mathrm{R} \lesssim 2.5 \times 10^8\,\mathrm{GeV} \quad \text{for} \quad m_{3/2} = 100\ \mathrm{GeV}. \tag{10}$$

Subsequent work [10] has shown that the true constraint is actually more restrictive than in ref.[27] by a factor upto ~ 2 for a radiative lifetime greater than $\sim 2 \times 10^7$ sec (corresponding to $m_{3/2} \lesssim 300\,\mathrm{GeV}$) but less stringent for shorter lifetimes. Here we have taken the lifetime for gravitino decay into a gauge boson-gaugino pair, $\tau_{3/2} \sim 4M_{\mathrm{Pl}}^2/N_{\mathrm{c}}m_{3/2}^3$ where N_{c} is the number of available channels (assuming $m_{\tilde{\gamma},\tilde{g}} \ll m_{3/2}$), so that [9]

$$\begin{aligned} \tau_{3/2\to\tilde{\gamma}\gamma} &\simeq 3.9 \times 10^5\,\mathrm{sec} \quad \left(\frac{m_{3/2}}{\mathrm{TeV}}\right)^{-3}, \\ \tau_{3/2\to\tilde{g}g} &\simeq 4.4 \times 10^4\,\mathrm{sec} \quad \left(\frac{m_{3/2}}{\mathrm{TeV}}\right)^{-3}. \end{aligned} \tag{11}$$

For a gravitino mass of 1 TeV, the radiative lifetime is about 4×10^5 sec and the relevant constraint now comes from requiring that the photofission of deuterium not reduce its abundance below the observational lower limit [28, 29]. A careful calculation shows that this requires $m_{3/2}(n_{3/2}/n_\gamma) \lesssim 5 \times 10^{-10}\,\mathrm{GeV}$ [10] so the reheat bound is weakened to $T_{\mathrm{R}} \lesssim 2 \times 10^9$ GeV. Photofission processes become ineffective for $\tau \lesssim 10^4$ sec but now there are new constraints from the effect of hadrons in the showers on the ^{4}He abundance. If the gravitino mass is 10 TeV with a corresponding lifetime of $\tau_{3/2\to\tilde{g}g} \sim 50$ sec, this constraint is $m_{3/2}(n_{3/2}/n_\gamma) \lesssim 1.5 \times 10^{-8}$ GeV [30], hence $T_{\mathrm{R}} \lesssim 6 \times 10^9$ GeV. There is *no* bound on T_{R} for gravitinos of mass exceeding 50 TeV which decay before nucleosynthesis [30, 8] but such a large mass cannot be accomodated in minimal models. Hence we take the reheat bound corresponding to the relic abundance in eq.(9) as given in ref.[10],

$$\begin{aligned} T_{\mathrm{R}} &\lesssim 10^8,\ 2 \times 10^9,\ 6 \times 10^9\ \ \mathrm{GeV}\,, \\ \text{for}\quad m_{3/2} &= 10^2, \qquad 10^3, \qquad 10^4\ \ \mathrm{GeV}\,. \end{aligned} \tag{12}$$

We emphasize that these are *conservative* bounds. For example in ref.[31] the relic gravitino abundance is calculated to be a factor of ~ 4 higher than in eq.(9) after including interaction terms between the gravitino and chiral multiplets; this would tighten all the above bounds on T_{R} by the same factor. However, the authors of ref.[31] also numerically calculate the cosmological constraint based on D+^{3}He photoproduction to be, inexplicably, a factor of ~ 30 more stringent than the one calculated analytically in ref.[10], which was confirmed by the full Monte Carlo simulation of ref.[32]. Consequently they obtain bounds on the reheat temperature which are much more severe than in eq.(12), for example $T_{\mathrm{R}} \lesssim 2 \times 10^6$ GeV for $m_{3/2} \sim 100$ GeV.

Recently the thermal production of gravitinos has been reexamined [33]; it is argued, contrary to previous studies [9], that massive gravitinos *can* achieve equilibrium with the background plasma at temperatures well below the Planck scale via interactions of their longitudinal spin-1/2 (goldstino) component with a cross-section which increases as T^2 due to the breaking of SUSY by finite temperature effects. The relic abundance is then given by [33]

$$\frac{n_{3/2}}{n_\gamma} \approx \frac{g_*^{1/2}\alpha_{\mathrm{s}}^3 T^3}{m_{3/2}^3 M_{\mathrm{Pl}}^2} \quad \sim 3 \times 10^{-13} \left(\frac{T_{\mathrm{R}}}{10^5\,\mathrm{GeV}}\right)^3 \left(\frac{m_{3/2}}{\mathrm{TeV}}\right)^{-2}, \tag{13}$$

where g_*, the number of relativistic degrees of freedom, equals 915/4 in the minimal supersymmetric standard model (MSSM) at $T \gg m_Z$. Imposing the nucleosynthesis

constraints discussed above then yields a very restrictive upper bound on the reheat temperature:

$$
\begin{array}{rcl}
T_{\rm R} & \lesssim & 2\times 10^4,\ \ 10^6,\ \ 2\times 10^7\ \ {\rm GeV}\ , \\
{\rm for}\quad m_{3/2} & = & 10^2, \qquad 10^3, \qquad 10^4\ \ {\rm GeV}\ .
\end{array} \tag{14}
$$

However this claim has been questioned on general grounds [34] and an explicit calculation [35] of finite-temperature effects demonstrates that these do not alter the estimate in eq.(9). Hence we will consider the impact on inflationary models of the reliable bound (12) but also note the implication of the proposed new bound (14), assuming a nominal gravitino mass of 1 TeV.

Another aspect of the 'gravitino problem' is the need to avoid direct production of gravitinos from the decay of the inflaton field during the reheating process [6], given that there is always a gravitational strength coupling between the gravitino (or any other state) and the inflaton. Unlike thermal production discussed above, this needs to be considered in the context of a specific model and we shall do so in Section 3.1.

2.2.2 Normalization to COBE

The semi-classical equation of motion for the inflaton field is [2]

$$\ddot{\phi} + 3H\dot{\phi} + V'(\phi) = 0. \tag{15}$$

Essentially all models generating an exponential increase of the scale-factor a satisfy the slow-roll conditions [36]

$$\dot{\phi} \simeq -\frac{V'}{3H}\ , \quad \epsilon \equiv \frac{M^2}{2}\left(\frac{V'}{V}\right)^2 \ll 1, \quad |\eta| \equiv \left|M^2\frac{V''}{V}\right| \ll 1. \tag{16}$$

Inflation ends (i.e. $\ddot{a}$ becomes zero) when

$$\max(\epsilon, |\eta|) \simeq 1. \tag{17}$$

Given these conditions, the spectrum of adiabatic density perturbations is [37]

$$\delta_{\rm H}^2(k) = \frac{1}{150\pi^2}\frac{V_\star}{M^4}\frac{1}{\epsilon_\star}, \tag{18}$$

where $\star$ denotes the epoch at which a scale of wavenumber k crosses the 'horizon' H^{-1} during inflation, i.e. when $aH = k$. (Note that for a scale-invariant spectrum $\delta_{\rm H}$ equals $\sqrt{4\pi}$ times the parameter $\epsilon_{\rm H}$, cf. ref.[36].) The CMB anisotropy measured by COBE allows a determination of the perturbation amplitude at the scale k_1^{-1} corresponding roughly to the size of the presently observable universe, $H_0^{-1} \simeq\ 3000\ h^{-1}\,{\rm Mpc}$, where $h \equiv\ H_0/100\ {\rm km\ sec^{-1}\ Mpc^{-1}}$ is the present Hubble parameter. The number of e-folds from the end of inflation when this scale crosses the horizon is [36]

$$
\begin{aligned}
N_1 \equiv N_\star(k_1) \ \ \simeq 51 \ \ &+ \ln\left(\frac{k_1^{-1}}{3000h^{-1}\,{\rm Mpc}}\right) + \ln\left(\frac{V_\star}{3\times 10^{14}\,{\rm GeV}}\right) + \ln\left(\frac{V_\star}{V_{\rm end}}\right) \\
&- \frac{1}{3}\ln\left(\frac{V_{\rm end}}{3\times 10^{14}\,{\rm GeV}}\right) + \frac{1}{3}\ln\left(\frac{T_{\rm R}}{10^5\,{\rm GeV}}\right),
\end{aligned} \tag{19}
$$

where we have indicated the numerical values we anticipate for the various scales. The exact procedure for normalization is somewhat subtle [38] since it depends on the precise shape of the density perturbation spectrum and on whether there are other contributions to the CMB anisotropy such as from gravitational waves. The higher multipoles in the CMB anisotropy also depend on the assumed composition of the dark matter (i.e. 'cold' or 'hot') since this determines how the primordial spectrum is modified on scales smaller than the horizon at the epoch of matter-radiation equality i.e. for $k^{-1} < k_{\rm eq}^{-1} \simeq 80\ h^{-1}$ Mpc, assuming the dark matter to contribute the critical density [39]. However, these are small effects and well within the observational uncertainties so initially we assume that the anisotropy in the COBE data arises purely due to the Sachs-Wolfe effect on super-horizon scales ($k^{-1} > k_{\rm dec}^{-1} \simeq\ 180\ h^{-1}\,$Mpc) at CMB decoupling. The best fitting quadrupole moment obtained from the angular power spectrum of the 2-year data, $Q_{\rm rms-PS} \simeq\ 20\ \mu$K [40], then corresponds to

$$\delta_{\rm H} = \sqrt{\frac{48}{5}}\frac{\langle Q\rangle}{T_0} \simeq 2.3 \times 10^{-5} \tag{20}$$

with an uncertainty of $\sim$ 10%, taking $T_0 \simeq 2.73$ K. (This is about 30% higher than the value obtained previously [36] from the 1-year COBE data [3].) Using eq.(18) then gives

$$V_1^{1/4} \simeq 7.3\ \epsilon_1^{1/4} \times 10^{16}\,{\rm GeV}, \tag{21}$$

demonstrating how the inflationary scale is strictly bounded from above [36], thus justifying our neglect of gravitational corrections to the potential. Applied to the supergravity potential following from eqs.(5) and (8) we have

$$V_1 = \Delta^4\ {\rm e}^{|\phi_1|^2/M^2}\left[\left|f'\left(\frac{\phi_1}{M}\right) + \frac{\phi_1^* f}{M}\right|^2 - 3|f|^2\right]. \tag{22}$$

For convenience we have chosen to define f in eq.(5) such that the RHS is just Δ^4, hence the COBE normalization requires

$$\frac{\Delta}{M} \simeq 3 \times 10^{-2}\epsilon_1^{1/4}. \tag{23}$$

At the end of inflation, the field ϕ begins to oscillate about its minimum until it decays, thus reheating the universe. In the class of models considered here, ϕ is a gauge singlet field in a hidden sector with only gravitational strength couplings to other states. Such models have been widely explored for they can readily produce an acceptable inflationary potential; clearly this class of models also has the best chance of satisfying the bound on the reheat temperature since here the inflaton has the smallest possible coupling to other states. The dominant coupling of ϕ to states χ in another sector with superpotential $P(\chi)$ has the form $(\partial V/\partial\phi)\ \ P(\chi)_A M^{-2}$, where the subscript A denotes that the chiral superfields in P should be replaced by their scalar components. We see that this generates a trilinear coupling to the light matter fields χ of strength $\sim \Delta^2/M^2$, corresponding to a decay width $\Gamma_\phi \sim [m_\phi/(2\pi)^3](\Delta^2/M^2)^2$. With our simplifying assumption that there are no small parameters in $f(\phi)$ the mass of the inflaton is

$$m_\phi \sim \frac{\Delta^2}{M}\ . \tag{24}$$

The inflaton thus decays at

$$t_{\rm R} \sim \Gamma_\phi^{-1} \sim (2\pi)^3 \frac{M^5}{\Delta^6} \tag{25}$$

and converts its energy content into radiation according to

$$\rho_\phi(t_{\rm R}) \simeq \frac{\pi^2}{30} g_*(T_{\rm R}) T_{\rm R}^4 \ . \tag{26}$$

(The ‘parametric resonance’ effect discussed in ref.[41] is irrelevant here since the inflaton has no coupling of the form $\phi^2\chi^2$ but only terms involving χ^3. This is because the supergravity couplings involve P^2 where P is the superpotential which is *trilinear* in the scalar fields, hence bilinear terms are suppressed by $m_\chi/M_{\rm Pl}$.) The temperature at the begining of the standard radiation-dominated era is thus

$$T_{\rm R} \sim \left(\frac{30}{\pi^2 g_*}\right)^{1/4} (\Gamma_\phi M)^{1/2} \simeq 2.2 \times 10^{-2} \frac{\Delta^3}{M^2} \ . \tag{27}$$

Demanding that this be less than the phenomenological bounds (eqs.12,14) then requires,

$$\begin{aligned} \left\{\frac{\Delta}{M}, \epsilon_1\right\} \quad &\lesssim \{2.7 \times 10^{-4},\ 6.1 \times 10^{-9}\} \qquad \text{for} \quad T_{\rm R} \lesssim 10^6\,\text{GeV}, \\ &\lesssim \{3.3 \times 10^{-3},\ 1.5 \times 10^{-4}\} \qquad \text{for} \quad T_{\rm R} \lesssim 2 \times 10^9\,\text{GeV}, \end{aligned} \tag{28}$$

where we have used eq.(23). Hence the constraint following from the need to inhibit thermal production of gravitinos forces the inflationary scale to be low; the COBE normalization then requires the scalar potential to be very flat indeed. In turn this implies that the ratio of the tensor to the scalar contributions to the observed CMB anisotropy, given by $R \simeq\ 12.4\epsilon_1$ [36], is negligible small. For example, taking the highest allowed value $\epsilon_1 \lesssim 1.5 \times 10^{-4}$ corresponding to the conservative reheat bound $T_{\rm R} \lesssim 2 \times 10^9\,$GeV, we have the constraint

$$R \lesssim 1.9 \times 10^{-3}. \tag{29}$$

This prediction can probably be observationally tested [42] by ongoing and forthcoming medium- and small-scale angular anisotropy experiments [43].

We have established the general conditions (28) which ensure a sufficiently low reheat temperature. Can these be achieved in realistic models? As discussed in Section 2.1.1, it is reasonable that a stage of symmetry breaking below the Planck scale should generate an inflationary potential with a related scale. Our example of gauge symmetry breaking gave $\Delta \sim 3 \times 10^{-5} M$ (eq.4), so one can easily obtain the required scale of inflation without any fine tuning. To discuss the small value required for the slope of the potential, we will have to specify what the dynamics of our inflationary scheme is, viz. new or chaotic. In new inflation [18], the inflaton rolls from an initial value near the origin towards its global minimum at the Planck scale; the derivative of the potential is neccessarily small in order to ensure sufficient inflation, hence such models *naturally* ensure a small value for ϵ_1. However in chaotic inflation [19], the inflaton starts rolling towards its minimum at the origin from an initially large value beyond the Planck scale. Hence for a generic power-law potential $V \propto \phi^\alpha$, MV'/V can be small only if the value of ϕ is much larger

than the Planck scale. The field value corresponding to the spatial scales probed by COBE is $\phi_1 \simeq (2\alpha\ N_1)^{1/2}\ M$, so the slope at this point is

$$\epsilon_1 \simeq \frac{\alpha}{4N_1}\ . \tag{30}$$

This is $\sim 2 \times 10^{-2}$ for $V \propto \phi^4$, hence violates even the conservative bound in eq.(28) by a factor of ~ 100. We conclude that it is not possible to satisfy the phenomenological constraints on the reheat temperature in any (one-scale) chaotic supergravity inflation model. (Conversely, such models predict a substantial tensor component in the CMB anisotropy, e.g. $R \simeq 25\%$ for a quartic potential.) We emphasize that by "chaotic" we specifically mean here models where the scalar field vev *decreases* during inflation from an initially large value beyond the Planck scale. We are not referring to the random initial conditions of chaotic models, which we do indeed adopt ourselves.

So far our analysis has been rather general. We now consider a specific example capable of giving inflation with small ϕ/M, which naturally satisfies the bound on the reheat temperature when the scalar density perturbations are normalized to the COBE observations. We will be able to study the remaining questions related to the reheat process in the context of this hybrid model which combines the dynamical evolution of new inflation with the initial conditions of chaotic inflation.[8]

3 A simple model

In this model [6] the inflation sector I, the SUSY breaking sector S and the visible sector G interact with each other only gravitationally and hence may be constructed separately in the superpotential. First consider the inflation sector. Requiring that SUSY remain unbroken in the global minimum, i.e.

$$\left|\frac{\partial I}{\partial \Phi} + \frac{\Phi^* I}{M^2}\right|_{\Phi=\Phi_0} = 0\ , \tag{31}$$

and setting the present cosmological constant to be zero,

$$V_I(\Phi_0) = 0\ , \tag{32}$$

implies

$$I(\Phi_0) = \frac{\partial I}{\partial \Phi}(\Phi_0) = 0\ . \tag{33}$$

Thus I can be expanded as a Taylor series about its minimum. The *simplest* form for I which satifies the above conditions is [6]

$$I = \frac{\Delta^2}{M}(\Phi - \Phi_0)^2\ , \tag{34}$$

where Δ is a mass parameter setting the energy scale for inflation. Now in order for successful inflation to occur by the slow roll-over mechanism, the scalar potential must be flat at the origin,

$$\frac{\partial V_I}{\partial \Phi}|_{\Phi=0} = 0\ , \tag{35}$$

[8] Such a hybrid model was first studied in ref.[20].

which sets $\Phi_0 = M$. This in turn sets

$$\frac{\partial^2 V_I}{\partial \Phi^2}|_{\Phi=0} = 0 \ , \tag{36}$$

since I does not contain cubic terms. (The flatness is a gauge-invariant property of the potential as Φ is a gauge-singlet; indeed Φ cannot carry any quantum numbers because it appears linearly in I.) The scalar potential obtained from eq. (8) is shown in Figure 1. The complex direction is stable while along the real direction we can expand

$$V_I(\phi) = \Delta^4 \left[1 - 4\left(\frac{\phi}{M}\right)^3 + \frac{13}{2}\left(\frac{\phi}{M}\right)^4 - 8\left(\frac{\phi}{M}\right)^5 + \frac{23}{3}\left(\frac{\phi}{M}\right)^6 + \ldots \right] . \tag{37}$$

Can such a form arise naturally? We noted earlier (see eq.5) that it was natural for the superpotential to acquire an overall mass scale of this form due to the underlying symmetries of the theory. We also argued that the initial value of the field, ϕ_0, corresponding to the point at which the first derivative vanished would dominate the final state of the universe. Expanding $V(\Phi/M)$ about this value,

$$V(\Phi/M) = a\,(m) + b\,(m)(\Phi - \Phi_i) + c\,(m)(\Phi - \Phi_i)^2 + \ldots \ , \tag{38}$$

where the coefficients depend on the moduli m. The coefficient $a(m)$ determines the value of the potential initially and hence the moduli will flow to minimize this. However if the other coefficients depend on independent combinations of the moduli they will be undetermined at this stage as they do not affect the initial vacuum energy. It is usually implicitly assumed that this is *not* the case, since all the moduli are supposed to get SUSY-breaking masses from gravitational coupling to the vacuum energy during inflation. However we stress that this is not likely for the random initial conditions assumed here. For example the simple Kähler potential $K = \sum_j \Phi_j^\dagger \Phi_j$ leads to the potential

$$V_I = \mathrm{e}^{\sum_j |\Phi_j|^2/M^2} \left[\sum_k \left| \frac{\partial I}{\partial \phi_k} + \frac{\phi_k^* I}{M^2} \right|^2 - \frac{3|I|^2}{M^2} \right] . \tag{39}$$

During inflation one term in the square brackets must be non-zero and thus one expects all fields ϕ_i, including moduli fields, to get a positive mass squared term. Minimizing this potential will fix all the moduli vevs independently during inflation. However this form, if it applies at all, will do so only about an enhanced symmetry point corresponding to vanishing vevs for the Φ_i. For initial conditions far from this point, K (and the corresponding exponential in eq.39) will be a function of $\Phi_j \langle \Phi_j^* \rangle + \Phi_j^* \langle \Phi_j \rangle$ and $\Phi_j^\dagger \Phi_j$, so minimization will lead to a correlated relation between the moduli fields, leaving other independent combinations of moduli unconstrained. As discussed earlier, we then expect that the initial conditions will allow *some* region in which the value of $b\,(m)$ is just that needed to make the second derivative of the potential vanish and this will dominate the final state of the universe because of the enhanced amount of inflation it will undergo. Thus we see that most of the features incorporated in eq.(37) following from eq.(34) are indeed natural and to be expected in any theory which has a potential with a turning point. The exceptional property is that I is a perfect square. Recall that this was

required to ensure the vanishing of the cosmological constant at the minimum (see eq.32). By adjusting the term $c\,(m)$ in eq.(38) we can always arrange that V vanishes, but not in a natural way.

However it is not difficult to construct an alternative model which does not suffer from this problem at all. All that is needed is a reason why, at the end of inflation, the superpotential is at least quadratic in the inflaton field expanded about the minimum of the potential. This suggests an underlying symmetry, Z_2 or larger. The example above does not have this symmetry because the Kähler potential is not quadratic about $\Phi = M$. It is straightforward to construct a variant which does have such a Z_2 symmetry via the superpotential $P = \Delta^2 M(\Phi^2 + a\Phi^4 + \mathcal{O}(\Phi^6)\ldots)$ and a Kähler potential $K = \Phi^\dagger\Phi$. Now clearly the associated potential vanishes at the minimum as required. In this case the coefficient a must be fine-tuned to give a potential with a turning point away from $\Phi = 0$ (which will now be the initial value) at which the second derivative also vanishes. However, as stressed above, in the case that a is determined by moduli fields, the fine tuned value will *automatically* be selected when considering the most likely state of the universe. We will not pursue this example further since the model already introduced contains the significant features of this type of inflationary model. Thus henceforth we will use the model of eqs.(34) and (37), which has this cancellation, as we proceed to consider the reheat phase.

Integrating eq.(15) back from the end of inflation, the field value corresponding to the spatial scales probed by COBE is $\phi_1 \simeq\ M/[12(N_1+2)]$ using $V'/V \simeq -12\phi^2/M^3$ from eq.(37), so

$$\epsilon_1 \simeq 72\left(\frac{\phi_1}{M}\right)^4 , \tag{40}$$

which is 4.4×10^{-10} for $N_1 = 51$, i.e. as small as is required to solve the gravitino problem. (Note that for the ubiquitous $\lambda\phi^4$ potential in chaotic inflationary scenarios, such a small slope can only be obtained for $\phi \gtrsim\ 10^5\ M$!) The total number of e-folds of inflation is

$$N_{\rm tot} \equiv \ln\frac{a(\phi_{\rm end})}{a(\phi_{\rm ini})} = \frac{1}{M^2}\int_{\phi_{\rm ini}}^{\phi_{\rm end}} -\frac{V}{V'}{\rm d}\phi \simeq \frac{M}{12\Delta}\ , \tag{41}$$

setting $\phi_{\rm ini} \sim \Delta$. Thus we need $\Delta/M \lesssim 1.7\times10^{-3}$ to get $N_{\rm tot}$ as large as is neccessary to solve the cosmological horizon/flatness problems. Since the potential (37) is dominated by a *cubic* term, the spectrum of scalar density perturbations will depart from scale-invariance at large k as $\sim \ln^2(Hk^{-1})_\star$ [44]. This corresponds to a 'tilted' spectrum with slope given by $n = 1 + 2\eta_1 - 6\epsilon_1$ [36], so for the potential (37), we get

$$n \simeq \frac{N_1 - 2}{N_1 + 2} \simeq 0.9\ , \tag{42}$$

which has less power on galactic scales in a CDM cosmogony and thus will provide a better match to observations [45].[9] Because of the tilt, the normalization to the COBE anisotropy is now somewhat higher corresponding to $\delta_{\rm H} \simeq 2.5\times10^{-5}$ [38] for the perturbation amplitude at large scales. This gives

$$\frac{\Delta}{M} \simeq 1.4\times10^{-4}, \tag{43}$$

[9]Note that the tilt of the spectrum is *not* associated with a significant tensor component in the CMB anisotropy (although the converse is true [42]), as is also the case with the 'natural' inflation model [46].

i.e. the inflationary scale is about 3×10^{14} GeV and the inflaton mass is about 5×10^{10} GeV. The reheat temperature is then

$$T_{\rm R} \sim 1.5 \times 10^5 \, \text{GeV}, \tag{44}$$

well below the conservative bound (12) and consistent with even the recently proposed bound (14).

3.1 Direct gravitino production

It is not sufficient to inhibit the thermal production of gravitinos; it is necessary, in addition, to suppress their production via inflaton decay. This is a particularly important for an inflaton in the hidden sector for all its couplings are gravitational and therefore comparable to all states. Thus the solution to the reheating constraint proves to be the problem where direct production is concerned. The couplings to the gravitino can indeed be suppressed but this is very model dependent. The best we can do here is to illustrate the problem and present a possible solution in the context of the model introduced above.

The couplings of the inflaton with interactions specified by the superpotential of eq.(34) to the gravitino is contained in the general coupling

$$\frac{|I_s|}{M^2}\bar{\psi}^{\mu}_{3/2}\sigma_{\mu\nu}\psi^{\nu}_{3/2} \ , \tag{45}$$

where I_s is the superpotential with superfields replaced by their scalar components. This leads to the coupling

$$h_{\phi\psi_{3/2}\psi_{3/2}} = 2\frac{\Delta^2(\phi - M)}{M^3} \ . \tag{46}$$

On the other hand the coupling of the inflaton to the scalar components of matter fields may be read from eq.(37). The dominant coupling comes from the interference term when the first term is expanded, giving

$$(\phi^* P_s + \phi P_s^*)\frac{2\Delta^2}{M^2} \ , \tag{47}$$

which contains quartic scalar couplings involving the inflaton and the matter fields in the full superpotential. In the MSSM the dominant coupling will be to the top squarks and the Higgs,

$$h_{\phi^*\tilde{t}\tilde{t}^c H_2} = h_t\frac{2\Delta^2}{M^2} \ , \tag{48}$$

where h_t is the top Yukawa coupling. We may use these couplings to deduce the relative production of gravitinos in inflaton decay. The crucial point is that although both couplings are gravitational in origin, there is a suppression factor $(\phi - M)$ in the gravitino coupling which follows because of the *square* appearing in the superpotential I. As discussed earlier, this is dictated by the empirical requirement that there be no cosmological constant after inflation and may be expected in any model in which the contribution to the cosmological constant from the inflationary sector vanishes. The inflaton decays at a time given by eq.(25), so using the virial theorem we may determine the mean value of the factor in eq.(46) to be

$$\langle(\phi - M)^2\rangle = \frac{\Delta^4}{16\pi^3 M^2} \ , \tag{49}$$

giving

$$\frac{\Gamma_{\psi_{3/2}\psi_{3/2}}}{\Gamma_{\tilde{t}\tilde{t}^c H_2}} \sim \frac{\Delta^4}{M^4} \ . \tag{50}$$

Is this suppression sufficient? The gravitinos produced by inflaton decay are highly relativistic since $m_\phi \sim 5 \times 10^{10}$ GeV. Until they become non-relativistic, their energy density relative to that of radiation remains constant. However after they do, and until they decay (or the universe becomes matter dominated) their relative energy density grows linearly with the expansion factor. This amounts to a growth of approximately 10^9 for a TeV mass gravitino. Thus the condition $\Delta/M \lesssim 6 \times 10^{-3}$ is adequate to make direct gravitino production phenomenologically acceptable. We see that this is amply satisfied by the value in eq.(43) fixed by the normalization to COBE.

3.2 The Polonyi problem

Lastly we consider whether a decoupled field will have stored potential energy during inflation which is released very late, thus recreating the problems associated with such a field that inflation was supposed to cure [11, 12]. The fields of concern abound in compactified (string) models [13]; they include the moduli mentioned above which are responsible for determining the couplings in the theory, the field related to the dilaton which determines the gauge coupling, and the moduli fixing the Kähler structure determining the shape and radius of compactification. The precise details of this problem are model dependent but we shall present a general description of the problem and determine under what conditions one may hope to avoid it. A detailed exposition of this question has recently been given in ref. [17]. While we are broadly in agreement, our emphasis on the likely solution is different and we have a new proposal as to how to deal with the moduli (as distinct from the dilaton). The question of weak scale moduli has also been recently discussed elsewhere [47, 48, 49, 50]; here we include a discussion of much lighter moduli of the type which arise if they are to determine the couplings of the Standard Model through their effect on the electroweak effective potential [51].

Suppose, as is expected in realistic models, that there is a stage of spontaneous symmetry breaking below the Planck scale and denote by δ the scale of the relevant potential. Now the moduli, m, which determine the couplings of the theory enter this potential only in the form of higher dimension operators suppressed by inverse powers of the Planck mass:

$$V(m) = \delta^4 \ k\left(\frac{m}{M}\right) \ . \tag{51}$$

where k is some function. If δ is the largest scale of spontaneous symmetry breaking then the vevs of the moduli are obtained by minimizing this potential. Of course it may be that this is not sufficient to determine all moduli, in which case one must consider further sectors of spontaneous breaking, as in models [51] where the couplings of the Standard Model are determined in this manner by the electroweak symmetry breaking potential.

The properties of the moduli field are very similar to those of the inflaton discussed above; its couplings are of gravitational strength and its mass is very small, suppressed by inverse powers of the Planck mass. Thus the moduli will have a lifetime given by

$$\Gamma_m \sim \frac{(m_m)^3}{M^2} \sim \frac{\delta^6}{M^5} \ . \tag{52}$$

We are concerned with the case $\delta < \Delta$ for then the moduli will decay, after inflation, at very late times and release too much entropy if the energy stored in the field is large. To estimate the magnitude of this stored energy we note that the moduli typically acquire a mass of $\mathcal{O}(\Delta^2/M)$ during inflation [12].[10] The sign of this term depends on the theory; if one takes the canonical kinetic term it will be positive and tend to stabilise the potential around $m = 0$ but if negative it will drive a large vev [53]. The expectations for the two cases are, respectively:

$$\begin{aligned} \langle m \rangle_+ &\sim M \left(\frac{\delta^4}{\Delta^4} \right) k'(0) \ , \\ \langle m \rangle_- &\sim M \left(\frac{\Delta^4}{\delta^4} \right) k''''(0) \ . \end{aligned} \tag{53}$$

In the second case, minimization of the simple form for the potential assumed above requires the moduli to be driven to large values, in excess of the Planck scale. This is probably unrealistic and all we will assume in this case is that the moduli are large, of $\mathcal{O}(M)$, but dependent on Δ. After inflation the vev of the moduli will clearly not depend on Δ. There are really only two possibilities, viz. $\langle m \rangle = 0$ or $\langle m \rangle \sim \mathcal{O}(M)$. The former can be shown to be a local minimum, for the moduli carry quantum numbers under discrete and/or other symmetries and the vanishing vev corresponds to a point of enhanced symmetry [54].

We can now quantify the problem. The energy stored in the potential during inflation and released afterwards as the moduli flow to their minima will be of $\mathcal{O}(\delta^4)$ *unless* the positions of the minima during and after inflation coincide to an accuracy much better than M. For the second possibility of eq.(53), it is unreasonable to suppose that the minima coincide to this accuracy. For the first possibility, the minima could coincide but *only* if after inflation the moduli vev remain zero corresponding to the enhanced symmetry point. Apart from this possibility, the energy in the moduli relative to the inflaton, just as inflation ends, is $R_m \sim (\delta/\Delta)^4$. However the important point to note is that the moduli are *not* in thermal equilibrium and will evolve according to their equation of motion analogous to eq.(15). From this it follows that the evolution of the field towards its minimum starts only at t_{end} when $H \sim \delta^2/M$. Until then, the energy in the moduli field is in the form of potential energy and thus grows relative to the energy in the inflaton field (or the products of the inflaton after reheating). At t_{end}, the moduli roll rapidly to their minimum, converting the potential energy into kinetic energy. However at this point $R_m \sim 1$, following immediately from the condition that the roll starts. Now both the inflaton and moduli are non-relativistic, so R_m remains constant until the time given in eq.(25) when the inflaton decays. Thereafter until the moduli decay (or the (non-moduli) universe becomes matter dominated), there will be a relative growth in the moduli energy proportional to the scale-factor. Following ref.[13] this can be estimated to be,

$$R_m = \min \left[\frac{\Delta^4}{\delta^4}, \left(\frac{\Delta^6}{10^{-56} M^4} \right)^{2/3} \right] \ . \tag{54}$$

[10] It is possible to suppress this at tree level in specific supergravity models [52] but typically the canonical order of magnitude will be restored in radiative order.

Clearly, any such growth is unacceptable as the universe will be matter dominated by the moduli, hence the entropy released by their subsequent decay will be far too large. We can envisage just three possible solutions to this problem:

(i) The moduli are fixed by a stage of symmetry breaking *before* inflation, i.e. $\delta > \Delta$. For example, new non-perturbative effects at a scale above that responsible for SUSY breaking may generate large moduli masses; however in all known examples these correspond to theories with negative cosmological constant, hence the universe necessarily suffers rapid gravitational collapse and does not evolve to a large enough size [17]. This does not exclude the idea altogether because there may well be string theories with such large non-perturbative effects and vanishing cosmological constant, but certainly one should find an example before concluding that this possibility is viable. Here we would like to point out a simple mechanism for generating moduli masses much larger than the SUSY breaking scale. We know that a large scale of symmetry breaking may be triggered by SUSY breaking along D- and F-flat directions. This comes about because a SUSY breaking scalar mass-squared for a field ρ which may be positive at the Planck scale can be driven negative at a lower scale, M_X, by radiative corrections. This will trigger symmetry breaking along a D- and F- flat direction (if there is one). The important point is that the vev induced is energetically favoured to lie close to M_X, *independent* of the SUSY breaking scale, even though it acts as the trigger. This vev in turn can generate moduli masses of $\mathcal{O}(M_X^2/M)$ from a term in the superpotential of the form $(mm'/M^2)\rho^3$, where the quadratic structure applies near the point of enhanced symmetry [54]. Thus we consider it quite likely that theories with a large intermediate scale of symmetry breaking will have moduli masses sufficiently large to avoid the Polonyi problem. However this mechanism does not work for the dilaton whose couplings have a different character. Although it is unlikely that the dilaton should acquire a mass much higher than the electroweak scale [17], it is not proven that this cannot happen through non-perturbative effects; indeed it has been conjectured that this does indeed happen [55]. In our opinion this would be the best solution, since otherwise it appears impossible to construct an inflationary potential at all. This is because in superstring theories the potential necessarily has a dependence on the dilaton and, with a light dilaton, the curvature of the potential in the dilaton direction is too large to permit inflation [21].

(ii) The moduli minima are the *same* during and after inflation. As stressed above, this is quite reasonable but only if this corresponds to a point of enhanced symmetry. The example given above illustrates how this comes about. However this explanation will not allow scenarios in which the moduli and associated Yukawa or soft SUSY breaking masses are determined at a low scale such as the electroweak breaking scale [51], because in this case the moduli vevs adjust in response to low energy phenomena (such as the alignment of soft masses with the fermion masses) which arise only on electroweak breaking. It seems exceedingly unlikely that the low energy minimum would, in this case, correspond to the minimum of the moduli during inflation.

(iii) The cosmology is non-standard after t_{R}. It has been suggested that there may have

been a brief secondary period of inflation (before baryogenesis) during which the moduli roll to their minima [48]; the following reheat phase then dilutes the moduli energy. (This is effectively equivalent to the possibility $\delta > \Delta$ but with a smaller Δ.) To be viable, this scenario must be carefully constructed — the second epoch of inflation must be short enough not to erase the density perturbations produced in the initial inflationary era, but long enough to adequately dilute the moduli energy. Now if the moduli are to roll to their minimum, the Hubble parameter during the second phase of inflation must be less than the moduli mass. For the case of moduli with electroweak scale masses considered in refs.[48, 49] this requires the vacuum energy during inflation to be at an intermediate scale of $M_{\rm I}^4 \sim m_W^2 M^2$. There should be fewer than $\sim 25 - 30$ e-folds of inflation if the density perturbations generated in the first stage of inflation are not to be erased. Finally the reheat phase must proceed through unsuppressed renormalizable couplings to Standard Model states in order to avoid reintroducing the moduli problem via the new inflaton. A particularly promising scenario for this second stage of inflation has been suggested recently [49]. It is noted that thermal effects can keep the field value at the origin from the temperature, $T_1 \sim (m_W M_{\rm I})^{1/2}$, at which the potential dominates, until the temperature, $T_2 \sim m_W$, at which the supersymmetry breaking mass triggers the phase transition. Thus the number of e-folds of inflation is automatically limited. Further it is plausibly argued that the relatively low scale of symmetry breaking needed to solve the moduli problem (i.e. $M_{\rm I} \sim 10^9\,{\rm GeV}$) is expected in supersymmetric theories with flat directions. Such an epoch of 'thermal' inflation is expected to occur irrespective of previous stages of inflation. Thus there is no conflict with our criterion of naturalness, i.e. we do not require multiple *correlated* periods of inflation. However, while this may cure the problem for moduli with weak scale masses, it cannot do so for much lighter moduli because the energy stored in such fields is not released until the Hubble parameter falls below their mass and they start to roll. In particular the moduli whose vevs are determined at the electroweak scale have masses of $\mathcal{O}(m_W^2/M)$, and their energy will *not* be diluted by inflation at an intermediate scale. To invoke the solution of refs.[48, 49] one now needs inflation at the weak scale with an inflationary potential of $\mathcal{O}(m_W^4)$. However the reheat temperature is then far below the value of $\sim 10\,{\rm MeV}$ needed to avoid disrupting standard nucleosynthesis. One can argue that along a flat direction the potential may be of $\mathcal{O}(m_W^4)$ while still having renormalizable couplings for the inflaton; however, following the argument presented in ref. [48], the reheat temperature will still be rather low, $T_{\rm R} \sim g^{2/3} m_\phi^{5/6} M^{1/6}$, where g is a gauge or Yukawa coupling. Taking $m_\phi \approx m_W^2/M$, this yields a maximum reheat temperature of $\mathcal{O}(10^{-6})\,{\rm MeV}$! Given this problem we know of no way to reconcile moduli of mass $\mathcal{O}(m_W^2/M)$ with the cosmological constraints discussed above if their vevs are determined at the electroweak breaking scale.

4 Conclusions

We have re-examined the possibilities for and the problems associated with supersymmetric inflation of the type that one might expect in compactified string theories. In the

effective supergravity theory following from such theories below the Planck scale, there are many hidden sectors corresponding to fields with only gravitational strength couplings. Such sectors offer very plausible inflaton candidates because the fields have rather small masses, related to spontaneous symmetry breaking below the compactification scale. As a result, inflation, if it occurs, will be associated with a scale far below the Planck scale, just as is required if an acceptable spectrum of density perturbations is to be generated. Whether there is an inflationary era at all in such theories depends on the initial conditions and the form of the effective potential. We have argued that weakly coupled fields which are not in thermal equilibrium are likely to have random initial conditions. As a result the final state of the universe is dominated by regions which have undergone an inflationary era, corresponding to a choice of initial conditions of the moduli for which the potential is anomalously flat and to an initial value of the inflaton near the flat part of its potential. Thus we consider that all the ingredients of an acceptable inflationary potential are present in effective supergravity theories without the need to fine tune the parameters or to seek forms of the potential such as in the chaotic inflation scenario which inflate for a wide range of initial conditions. An analysis of the possible origin of the scale of inflation suggests that it should be bounded above by the scale of SUSY breaking. If this arises through gaugino condensation then the relevant scale may be of $\mathcal{O}(10^{13})$ GeV. Inflationary models with a flat potential and sub-Planck scale vevs require such an intermediate scale to generate the correct level of density perturbations. On the other hand chaotic inflationary models require a much larger scale.

The main difficulty encountered by supergravity inflationary schemes concerns the cosmology *after* inflation. This is because the many hidden sector fields which were desirable to give candidate inflatons now cause severe problems. Due to their weak coupling they decay very late in the evolution of the universe and, if they contain a sizeable fraction of the total energy density, produce unacceptable amounts of entropy in the process. We discussed three aspects of this problem in some detail. The first is the thermal production of such weakly coupled states, e.g. the gravitino, due to reheating after inflation. We noted that the hidden sector inflaton avoids this problem in a very natural way for, having only gravitational strength couplings, it reheats the universe to a low temperature. In simple models with only a single stage of inflation it is possible to relate the reheat temperature directly to the spectrum of density perturbations. We found that normalizing the latter to the COBE data constrains the reheat temperature to be $\sim 10^5$ GeV, comfortably within the phenomenological bounds for a TeV scale gravitino. An important, and observationally testable, implication is that there should be a negligible tensor component in the CMB anisotropy and that the power spectrum of scalar density perturbations should be tilted away from scale-invariance. The models which provide this nice consistency have small values of the inflaton field relative to the Planck scale. In contrast, models of the chaotic type (e.g. ref.[56]) give too high a reheat temperature if they are normalized to generate the correct amplitude of density perturbations. Of course it is possible to evade these conclusions by postulating non-standard evolution after the inflationary era but such schemes may require a measure of fine tuning if they are not to destroy the density perturbations produced by the initial stage of inflation.[11] Moreover, given the success of the simplest schemes, it seems to us reasonable to consider their predictions

[11] As just discussed, an exception is thermal inflation [49].

as the paradigm, rather than those of more involved models, e.g. involving two coupled fields [57], which are constructed to avoid problems not encountered in the simple ones.

The second potential problem for the supergravity inflation models is the direct production of gravitinos and other weakly coupled states through inflaton decay. This is more model dependent than thermal production and we have considered this problem only for minimal supergravity. We find that the *empirical* requirement that there be vanishing cosmological constant in the inflaton sector after inflation provides the required suppression of the production of gravitinos relative to matter fields. Again we find that the bounds for a TeV mass gravitino are easily satisfied for a low inflationary scale.

Finally we considered the Polonyi problem, viz. the difficulty in suppressing the potential energy stored during inflation in weakly coupled hidden sector fields such as the moduli of compactified string theories. We considered three possible solutions. The first is that all moduli have vevs fixed at a scale above inflation. The second is that the minimum of the potential during inflation coincides with the minimum after inflation. The third is a late stage of inflation. In all cases the implication is that the moduli *cannot* be treated as dynamical variables at the electroweak scale determining the couplings in the low energy theory.

The low reheat temperature requires the baryon asymmetry of the universe to be generated at a relatively low energy scale and there is a natural mechanism [58] to achieve this in the context of supergravity inflation. We are presently examining this question as well as the detailed expectations in our model for observations of large-scale structure [59].

Acknowledgement: We wish to thank Steven Abel, Maria Bento, Orfeu Bertolami, Pierre Binétruy, Savas Dimopoulos, Gian Giudice, Costas Kounnas, Fernando Quevedo, Lisa Randall, Ewan Stewart and Fabio Zwirner for helpful comments and criticism. We are particularly grateful to Andrei Linde and David Lyth for clarifying discussions.

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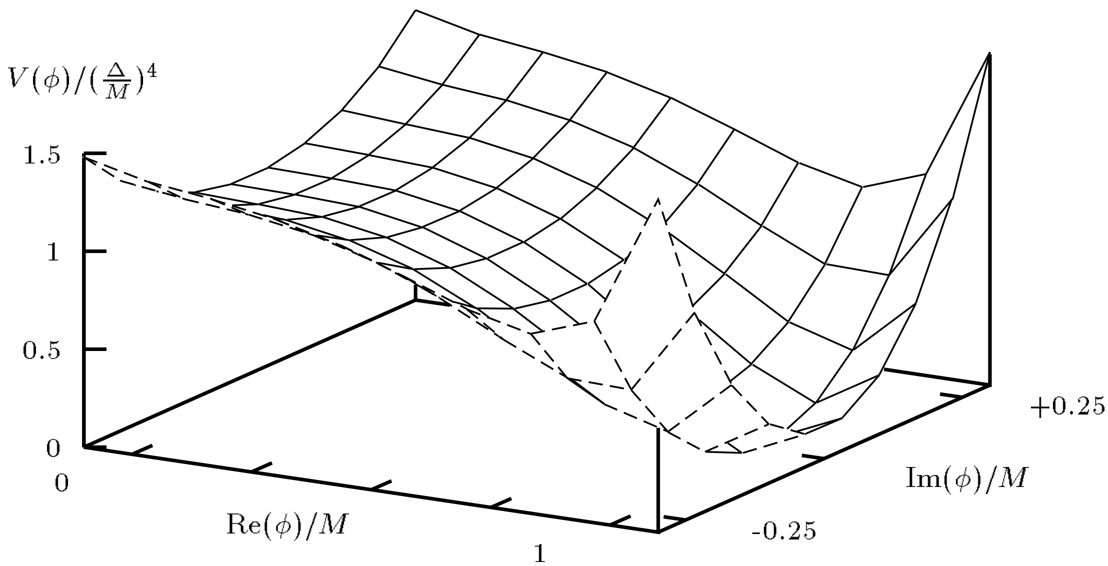

Figure 1: The complex scalar inflationary potential.